Same-sign single dilepton productions at the LHC

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Abstract

We examine the same-sign single dilepton productions of $\ell_i^{\pm}\ell_j^{\pm}$ ($\ell_{i,j} = e, \mu$) in high-energy proton-proton collisions at the Large Hadron Collider (LHC) in models with doubly charged Higgs scalars as well as heavy Majorana neutrinos. We demonstrate that these spectacular productions can be detected at the LHC for a class model in which the doubly charged Higgs scalars couple only to the right-handed charged leptons. The ranges of the possible doubly charged Higgs masses and mixings to observe the processes at the LHC are discussed.

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It is well accepted that neutrinos have very small masses. However, the origin of such smallness remains unclear. One of the most popular solutions is that they arise from the seesaw mechanism with one or more right-handed heavy Majorama neutrinos (HMN). On the other hand, without right-handed neutrinos, it is well known that the simplest way to have a Majorana neutrino mass term in the standard model (SM) is to introduce a complex triplet Higgs T with the hypercharge of $Y = -2$, defined by

$$T = \begin{pmatrix} T^0 & \frac{T^-}{\sqrt{2}} \\ \frac{T^-}{\sqrt{2}} & T^{--} \end{pmatrix}, \tag{1}$$

which can couple to $SU(2)_L$ lepton doublets (L_{iL}) [1]

$$\mathcal{L}_L = g_{ij}\overline{L^c_{iL}}T^\dagger L_{jL} + \text{H.c.}, \tag{2}$$

where g_{ij} are the coupling constants, $i, j = e, \mu, \tau$ and c stands for the charge conjugation. The neutrino masses are generated to be $g_{ij}v_T$ after the triplet scalar T receives the vacuum expectation value (VEV) of v_T. Since the major goals of the Large Hadron Collider (LHC) are searching for Higgs scalars and understanding the mechanism of the fermion mass generation, the HMNs and the triplet Higgs should be parts of the studies at the LHC.

The most interesting models which contain the triplet are left-right symmetric and little Higgs models [2, 3, 4]. Phenomenologically, the doubly charged scalar in the complex triplet could decay into the like-sign dileptons ($T^{\pm\pm} \to \ell_i^\pm \ell_j^\pm$) with a high invariant mass, which provides a spectacular signature with a relatively small background [5] at hadron colliders. A current limit set by the direct search at the Tevatron in Fermilab is $M_{T^{\pm\pm}} > 136$ GeV [6], in which the Drell-Yan (DY) annihilation processes $q\bar{q} \to \gamma^*, Z^* \to T^{++}T^{--}$ to the final states of $e^\pm e^\pm, e^\pm\mu^\pm, \mu^\pm\mu^\pm$ were used and a long-lived doubly-charged scalar, corresponding to $g_{ij} \gtrsim 10^{-5}$, was assumed. However, from the current neutrino mass upper bounds [7] of 0.1 eV, extracted from the neutrino oscillation data and cosmological experiments, the coupling of g_{ij} cannot be large if v_T is not too small. Note that $v_T \leq 4.41$ GeV [8, 9] is constrained by the precision data of $\rho = 1.002^{+.0007}_{-.0009}$ [7].

Accordingly, one concludes that in the model with $\mathcal{L}_L$ of Eq. (2), the production of $T^{\pm\pm}$ in the W-boson fusion decaying into a like-sign single dilepton, $W^\pm W^\pm \to T^{\pm\pm} \to \ell_i^\pm \ell_j^\pm$, are too small to be found generally at the LHC due to the following reasons: (a) the production rates are proportional to $(v_T/v)^2$ which is numerically small even v_T is set to be close to

the upper limit; and (b) as $g_{ij} \lesssim 10^{-10}$ for $v_T \sim 4.41$ GeV, the widths of $T^{\pm\pm} \to \ell_i^\pm \ell_j^\pm$ are very small and other channels would be opened to dominate over these dileptons signatures [10]. Moreover, small coupling constants g_{ij} are needed in order to fit the neutrino mixing matrix.

Recently, a model was proposed [8] with the $SU(2)_L$ complex triplet $T_{(-2)}$ and an additional doubly-charged singlet $\Psi_{(4)}$ to the SM, where the subscript denotes the hypercharge. In the model, a new Yukawa interaction, involving $SU(2)_L$ charged lepton singlets (ℓ_R),

$$\mathcal{L}_R \;=\; Y_{ij}\overline{\ell^c_{iR}}\ell_{jR}\Psi + \text{H.c.}, \tag{3}$$

is introduced due to Ψ but the one in Eq. (2) is forbidden by imposing some symmetry for the Higgs fields such as

$$\phi \to +\phi,\ \phi' \to -\phi',\ T \to -T \text{ and } \Psi \to +\Psi, \tag{4}$$

where an extra Higgs doublet ϕ' has been also included. However, since the extra doublet leads to no new effects [11] on the fermion couplings, the structure of the doubly-charged Higgs scalars as well as the phenomenology in Refs. [8, 9], we will not discuss it further here. In this model, as the neutrino masses are generated radiatively at two-loop level [8, 9, 12], the small neutrino mass problem can be naturally understood even with $Y_{ij} = O(1)$ and v_T around the upper limit simultaneously.

In this paper, we concentrate on doubly-charged scalars of $T^{\pm\pm}$ and $\Psi^{\pm\pm}$. The two fields can form doubly-charged massive physical states $P_1^{\pm\pm}$ and $P_2^{\pm\pm}$ with the mixing angle δ. It was argued in Ref. [9] that at least one of the doubly charged Higgs scalars is well within the reach of the LHC. We take P_1 to be this (lighter) state and focus on its phenomenology. In particular, we investigate the processes $pp \to P_1^{\pm\pm}X \to \ell_i^\pm \ell_j^\pm X$ under the conditions of the LHC: $\sqrt{s}$=14 TeV and L=320 fb^{-1}, where $\sqrt{s}$ is the beam energy and L is the integrated luminosity per year. We will also consider the contributions to the processes due to the HMNs. We will choose the condition

$$\sigma L \geq n \tag{5}$$

as n events of the observation criteria for the process, where σ denotes the cross section.

We start by evaluating the differential cross sections for the processes

$$pp \to \ell_i^\pm \ell_j^\pm X \tag{6}$$

at the LHC via the intermediate doubly charged Higgs scalar $P_1^{\pm\pm}$ by neglecting the transverse polarizations of W bosons and quark mixings, where X represents 2 jets, denoted as JJ. The leading-order Feynman diagram for the processes in Eq. (6) is shown in Fig. 1. For $\ell_{i,j} = e$ or μ, one has spectacular signatures of the same-sign dilepton pairs of $e^\pm e^\pm$ or $\mu^\pm\mu^\pm$ or $e^\pm\mu^\pm$ without missing energy. For the modes with one or two τ leptons, the final states are the above dilepton pairs but with missing energy or pions with missing energy. In this study, we shall not discuss the productions with missing energy as they are suppressed. According to Ref. [9], the gauge-scalar and the lepton-scalar couplings are given by

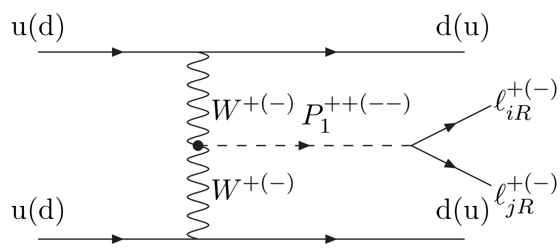

FIG. 1: Feynman diagram for $pp \to \ell_i^\pm \ell_j^\pm JJ$ mediated by $P_1^{\pm\pm}$.

$$\frac{g^2}{\sqrt{2}} \upsilon_T c_\delta W_\mu^+ W_\nu^+ P_1^{--} + \text{H.c. and } Y_{ij} s_\delta P_1^{--} \overline{\ell_{iR}^c} \ell_{jR} + \text{H.c.}, \tag{7}$$

respectively, where $c_\delta \equiv \cos\delta$ and $s_\delta \equiv \sin\delta$. The decay of $P_1^{\pm\pm}$ can proceed by four types of channels: $P_1^{\pm\pm} \to \ell_{iR}^\pm \ell_{jR}^\pm$, $P_1^{\pm\pm} \to W^\pm W^\pm$, $P_1^{\pm\pm} \to W^\pm P^\pm$ and $P_1^{\pm\pm} \to W^\pm W^\pm T_a^0$, where $P^\pm$ and T_a^0 are the single-charged and neutral components of the Higgs scalars in the model, respectively. The decay widths are given by [9]

$$\Gamma(\ell_{iR}^\pm \ell_{jR}^\pm) = (1+\delta_{ij}) \frac{|Y_{ij}|^2}{16\pi} s_\delta^2 M_{P_1}, \tag{8}$$

$$\Gamma(W^\pm W^\pm) = \frac{g^4 \upsilon_T^2 c_\delta^2}{16\pi M_{P_1}} \sqrt{1 - \frac{4M_W^2}{M_{P_1}^2}} \left(3 - \frac{M_{P_1}^2}{M_W^2} + \frac{M_{P_1}^2}{4M_W^2}\right), \tag{9}$$

$$\Gamma(W^\pm P^\pm) = \frac{g^2 c_\delta^2 M_{P_1}^3}{16\pi M_W^2} \lambda^{3/2} \left(1, \frac{M_W^2}{M_{P_1}^2}, \frac{M_P^2}{M_{P_1}^2}\right), \tag{10}$$

where $\lambda(x,y,z) = x^2 + y^2 + z^2 - 2xy - 2xz - 2yz$. The three-body decay modes are expected to be relatively suppressed by the phase space compared to the two-body ones. In Fig. 2, we show the decay widths with two extreme cases of the mixing angles. The like-sign dilepton decays provide clean and almost negligible SM background signatures. Moreover, the branching ratios depend on the Yukawa couplings g_{ij} (Y_{ij} in our case) which are strongly

linked with the different scenarios for the neutrino mass generation mechanisms [10, 13, 14].

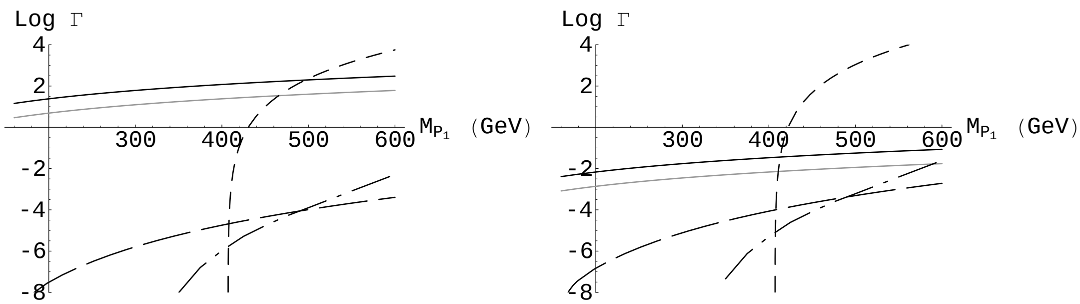

FIG. 2: Logarithms of the decay widths (in units of GeV) of $P_1^{\pm\pm}$ as functions of M_{P_1}, where the left (right) figure corresponds to the maximal (small) mixing of $\sin\delta = 1/\sqrt{2}$ (0.12), while the solid, dotted, long-dashed, short-dashed, and dot-dashed lines represent $\ell_i^\pm \ell_i^\pm$, $\ell_i^\pm \ell_j^\pm$ $(i \neq j)$, $W^\pm W^\pm$, $W^\pm P^\pm$, $W^\pm W^\pm T_a^0$ modes, respectively.

The differential cross sections for the processes in Eq. (6) are found to be

$$\frac{d\sigma_\pm^{pp}}{d\cos\theta} = A\left(\lambda_1^{ij}\right)^2 H_\pm^{pp}\,, \tag{11}$$

where θ is the angle between the directions for WW or qq and same sign leptons,

$$\begin{aligned} A &= \frac{G_F^4 M_W^6}{2^7\pi^5} = 50 \text{ ab},\ \ \lambda_1^{ij} = \sqrt{2-\delta_{ij}}\,|Y_{ij}|\,c_\delta s_\delta, \\ H_\pm^{pp} &= \left(\frac{v_T}{M_W}\right)^2 \int_{z_0}^1 \frac{dz}{z}\int_z^1 \frac{dy}{y}\int_y^1 \frac{dx}{x} p_\pm\left(x, xs\right) p_\pm\left(\frac{y}{x}, \frac{y}{x}s\right) l\left(\frac{z}{y}\right) h\left(\frac{s}{M_{P_1}^2}z\right), \end{aligned} \tag{12}$$

with $z_0 = M_{P_1}^2/s$. In Eq. (12), $h(t)$ are the normalized cross sections for the subprocesses of $W^\pm W^\pm \to \ell_i^\pm \ell_j^\pm$, given by

$$h\left(t\right) = \frac{t\left(t - 4M_W^2/M_{P_1}^2\right)}{\left(t-1\right)^2 + \Gamma_{P_1}^2/M_{P_1}^2}\,, \tag{13}$$

with the total decay width of $P_1^{\pm\pm}$ as:

$$\Gamma_{P_1} = 3\left[\Gamma(\ell_{iR}^\pm \ell_{iR}^\pm) + \Gamma(\ell_{iR}^\pm \ell_{jR}^\pm)_{i\neq j}\right] + \Gamma(W^\pm W^\pm) + \Gamma(W^\pm P^\pm) + \Gamma(W^\pm W^\pm T_a^0)\,, \tag{14}$$

$l(r)$ is the normalized luminosity (multiplied by r) of $W^\pm W^\pm$ pairs in the two-quark system [15], defined by

$$l\left(r\right) = -\left(1+r\right)\ln r - 2\left(1-r\right)\,, \tag{15}$$

and $p_\pm(x,Q^2)$ are the quark distributions in the proton, which have the forms:

$$p_+\left(x,Q^2\right) = x\sum_i q_i\left(x,Q^2\right) = x\left(u+c+t+\bar{d}+\bar{s}+\bar{b}\right), \tag{16}$$

$$p_-\left(x,Q^2\right) = x\left(\bar{u}+\bar{c}+\bar{t}+d+s+b\right). \tag{17}$$

It is interesting to note that the angular distributions in Eq. (11) are uniform on the quark level as there is only the s channel diagram for each of the processes in Eq. (6).

In the numerical calculation of the differential cross sections in Eq. (11), we use the CTEQ5 parton distributions [16] and take $|Y_{ij}|=1$, $s_\delta=0.12$ or $1/\sqrt{2}$, and v_T=4 GeV [9]. We note that $|Y_{ij}|=1$ is just a convenient choice. The constraints on these couplings from the neutrino oscillation data, rare decays, and $0\nu\beta\beta$ decays are studied in Ref. [9], where it is shown that the weakest upper bound is for the $\mu\mu$ production: $|Y_{\mu\mu}|<3.5$. However, due to $Y_{\ell\tau}<0.2$ $(\ell=e,\mu)$ and $|Y_{\tau\tau}|<0.02$ [9] as well as the small branching ratios of $\tau\to\ell\nu_\tau\bar{\nu}_\ell$, we shall exclude the productions with one or two taus in our discussion.

In Fig. 3, we show the relation between the cross sections and the mixing angle δ at $M_{P_1}=200$ GeV. Note that for $i\neq j$ there is an additional factor 2 in Eq. (11). In the figure, we also give the one event discovery limit (DL) at the LHC according to the observation criteria in Eq. (5). We find that since the decay widths depend on the mixing angle in different ways, the maximal cross section is not happened in the large mixing angle of $s_\delta=1/\sqrt{2}$ but around $s_\delta\sim 0.15$. In Fig. 4, we plot the cross sections of $P_1^{\pm\pm}$ and the DL at various values of M_{P_1}. The rate for P_1^{++} is about twice to that of P_1^{--} as expected based on the larger u-quark content in the proton at the LHC. In the case of $s_\delta=1/\sqrt{2}$, the processes via P_1^{--} are unobservable at the LHC. On the other hand, in the case of $s_\delta=0.12$, the cross sections drastically decrease as M_{P_1} is above 420 GeV because the decay channel of $W^\pm P^\pm$ opens up and becomes a dominant mode as seen from Fig. 2. We can conclude that, at the LHC, it is possible to detect the processes in Eq. (6) via the intermediate doubly-charged Higgs with its mass in the range from 180 GeV to 400 GeV while the mixing is between $\sin\delta=0.03$ and 0.85. It should be pointed out that one may tune the parameters to push the decay widths of $\Gamma(W^\pm P^\pm)$ at higher M_{P_1} to open up the modes of $W^\pm P^\pm$, then the searching range for $P_1^{\pm\pm}$ may be extended.

We now study the mechanism to produce dileptons in pp collisions due to the intermediate HMNs [17, 18]. The processes in Eq. (6) mediated by the HMN are illustrated in Fig. 5, where M_{N_1} corresponds to the mass of the lightest HMN. The differential cross sections are

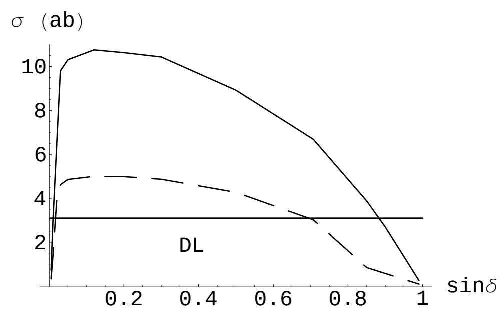

FIG. 3: Cross sections of $pp \to P_1^{++}X \to \ell_i^+\ell_j^+X$ (solid line) and $pp \to P_1^{--}X \to \ell_i^-\ell_j^-X$ (dashed line) as functions of $\sin\delta$ at $M_{P_1} = 200$ GeV, where the straight line is the one event discovery limit (DL) at the LHC.

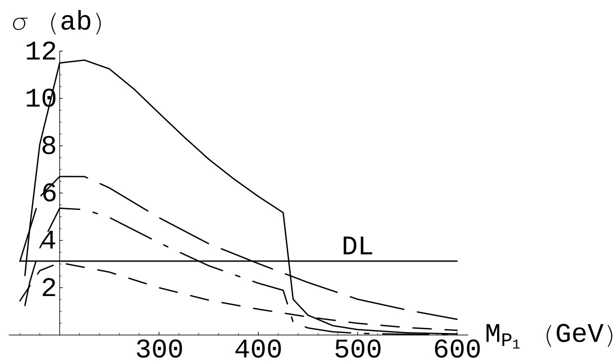

FIG. 4: Cross sections of $pp \to P_1^{\pm\pm}X \to \ell_i^\pm\ell_j^\pm X$ as functions of M_{P_1}, where the solid (dot-dashed) and long-dashed (short-dashed) lines stand for the processes $pp \to P_1^{++}X \to \ell_i^+\ell_j^+X$ $(\ell_i^-\ell_j^-)$ with $\sin\delta = 0.12$ and $1/\sqrt{2}$, respectively, while the straight line is the one event DL at the LHC.

given by

$$\frac{d\sigma_N^{pp}}{d\cos\theta} = 2A\left(\rho_1^{ij}\right)^2 N^{pp}\,, \tag{18}$$

with

$$\begin{aligned}
\rho_1^{ij} &= \sqrt{2-\delta_{ij}}\,|u_{i1}u_{j1}|\;, \\
N^{pp} &= \left(\frac{M_{N_1}}{M_W}\right)^2 \int_{\tilde{z}_0}^1 \frac{dz}{z}\int_z^1 \frac{dy}{y}\int_y^1 \frac{dx}{x} p_\pm(x, xs)\, p_\pm\left(\frac{y}{x}, \frac{y}{x}s\right) l\left(\frac{z}{y}\right) n\left(\frac{s}{M_{N_1}^2}z, \cos\theta\right)\,,
\end{aligned} \tag{19}$$

where $\tilde{z}_0 = 4M_{N_1}^2/s$, u_{i1} are the mixing matrix elements between the ith charged lepton and the heavy neutrino, and $n(t, \cos\theta)$ are the normalized cross sections for the subprocesses

$W^\pm W^\pm \to \ell_i^\pm \ell_j^\pm$, given by

$$n\left(t,\cos\theta\right) = \left(\frac{1-\cos\theta}{1-\cos\theta+2t^{-1}} + \frac{1+\cos\theta}{1+\cos\theta+2t^{-1}}\right)^2 . \tag{20}$$

Numerically, we find that these processes cannot be observed at the LHC even with a lower M_{N_1}. However, it is still possible at some higher luminosity colliders beyond the LHC [18]. It is interesting to note that the angular distributions in Eq. (18) given by the HMN mechanism are not uniform in contrast with the uniform ones in Eq. (11) by the doubly charged Higgs. Moreover, the produced same-sign leptons in Fig. 5 are left-handed, whereas those are right-handed in Fig. 1.

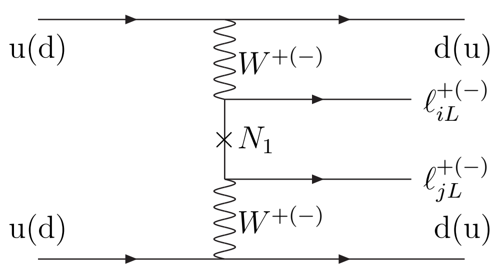

Fig. 5 Feynman diagram for the processes $pp \to \ell_i^\pm \ell_j^\pm JJ$ mediated by a heavy Majorana neutrino.

In summary, the possibility to observe the same-sign single dilepton productions at the LHC has been examined based on the doubly charged Higgs and Majorana neutrino mechanisms. We have demonstrated that the productions of $pp \to \ell_i^\pm \ell_j^\pm JJ$ $(\ell_{i,j} = e, \mu)$ can only be observed at the LHC in the model with the doubly charged scalars coupling to the right-handed charged leptons. In particular, the ranges of doubly charged Higgs masses and mixings to observe the productions in terms of the one-event discovery limits have been determined. We have also shown that the angular distributions of the differential cross sections for the processes are uniform on the quark level in contrast with the non-uniform ones due to the Majorana neutrino exchange mechanism. Finally, we remark that we have considered the processes $e^\mp p \to \ell_i^\mp \ell_j^\mp JJ$ via the intermediate $P_1^{\mp\mp}$ and HMN. However, by using the same method as for the pp collisions we find that these processes cannot be observed at the near future ep colliders [19].

Acknowledgements

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